How Hard Are the Form Factors in Hadronic Vertices with Heavy Mesons?

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Abstract

The $ND\Lambda_c$ and $ND^*\Lambda_c$ form factors are evaluated in a full QCD sum rule calculation. We study the double Borel sum rule for the three point function of one meson one nucleon and one Λ_c current up to order six in the operator product expansion. The double Borel transform is performed with respect to the nucleon and Λ_c momenta, and the form factor is evaluated as a function of the momentum Q^2 of the heavy meson. These form factors are relevant to evaluate the charmonium absorption cross section by hadrons. Our results are compatible with constant form factors in these vertices.
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QCD predicts that at high energy density, hadronic matter will turn into a plasma of deconfined quarks and gluons, the quark-gluon plasma (QGP), and there is a great deal of expectation in the comunity that the QGP will be observed in heavy ion collisions at the relativistic heavy ion collider (RHIC), which just started to operate at the Brookhaven National Laboratory. As a matter of fact, very recently the NA50 Collaboration reported [1] that the anomalous suppression of J/ψ observed in Pb+Pb collisions at CERN-SPS indicated already the formation of QGP, since these experimental data ruled out conventional hadronic models for the J/ψ suppression.

However, there are calculations that reproduce the new NA50 data reasonably well up to the highest transverse energies [2–4], based on hadronic J/ψ dissociation alone. Therefore, while there are suggestions that the anomalous suppression may be due to the formation of the QGP, other more conventional mechanisms based on J/ψ absorption by comovers and nucleons have to be still considered.

The main ingredient in the calculations based on hadronic J/ψ dissociation is the magnitude of the J/ψ absorption cross section by hadrons, which is not known experimentally. In refs. [2–4] this cross section was introduced as a free parameter. Therefore, in order to use the J/ψ suppression as a signature for the formation of the QGP in heavy ion collisions it is important to have a better knowledge on the interactions between charmonium states

and co-moving hadrons.

Various approaches have been used in evaluating the charmonium absorption cross section by hadrons. Some of them use meson exchange models based on hadronic effective lagrangians [5–9]. The couplings and form factors needed in these effective lagrangians are not phenomenologically known, and sometimes SU(4) relations are used to estimate them. However, we don't have any evidence that SU(4) relations should be taken seriously into account. Also, in general, it is assumed a monopole form factor at the hadronic vertex [7,8], with another unknown parameter (the cut-off). The results obtained for the cross section are very sensitive to the couplings and to the form factors [7–9]. As an example, in ref. [7] it was shown that the $\pi J/\psi \rightarrow DD^*$ cross section may vary by almost one order of magnitude depending on the value of the cut-off. Therefore, it is very important to estimate these couplings and form factors with a more theoretical approach.

In this work we use the QCD sum rule (QCDSR) method [10] based on the three-point function to evaluate the $ND\Lambda_c$ and $ND^*\Lambda_c$ form factors and coupling constants. In a previous work [11] we have already evaluated the $ND\Lambda_c$ coupling using the QCDSR. We have obtained $g_{ND\Lambda_c} = 6.7 \pm 2.1$ while SU(4) relations used in ref. [8] give $g_{ND\Lambda_c} = 14.8$. Of course the relevance of this difference can not be underestimated since the cross section is proportional to the square of the coupling constant. In this work we extend the calculation done before [11] and study the Q^2 dependence of the form factor for an off-shell heavy meson, which is the situation relevant for the calculations performed in refs. [6,8].

The three-point function associated with a $ND(D^*)\Lambda_c$ vertex is constructed with the two baryon currents, η_{Λ_c} and η_N, for Λ_c and the nucleon respectively, and the meson $D(D^*)$ current, $j_5(j_\mu)$ which we generically call j_M:

$$A_M(p,p',q) = \int d^4x d^4y \langle 0|T\{\eta_{\Lambda_c}(x) j_M(y) \overline{\eta}_N(0)\}|0\rangle e^{ip'x} e^{-iqy} \;, \tag{1}$$

with the currents given by [12,13]

$$\eta_{\Lambda_c} = \varepsilon_{abc}(u_a^T C\gamma_5 d_b) Q_c \;, \tag{2}$$

$$\eta_N = \varepsilon_{abc}(u_a^T C\gamma^\mu u_b)\gamma_5\gamma_\mu d_c \;, \tag{3}$$

$$j_5 = \overline{Q} i\gamma_5 u \;, \tag{4}$$

$$j_\mu = \overline{Q}\gamma_\mu u \;, \tag{5}$$

where Q, u and d are the charm, up and down quark fields respectively, C is the charge conjugation matrix and $q = p' - p$.

The general expression for the vertex function in Eq.(1) has four independent structures in the case of A_5 [11] and twelve in the case of A_μ [14]. In principle any of the four (twelve) invariant structures appearing in the A_M expression can be used to calculate the form factor and the sum rules should yield the same result. However, each sum rule could have uncertainties due to the truncation in the OPE side and different contributions from the continuum. Therefore, depending on the Dirac structure we can obtain different results due

to the uncertainties mentioned above. The tradicional way to control these uncertainties, and therefore to check the reliability of the sum rule, is to evaluate the stability of the result as a function of the Borel mass.

Recently, in ref. [15] it was pointed out that a better determination of $g_{\pi NN}$ can be done with the help of the $\gamma_5\sigma_{\mu\nu}$ structure, since this structure is independent of the effective models employed in the phenomenological side and it gets a smaller contribution from the single pole term coming from $N \rightarrow N^*$ transition. This was confirmed also in the case of the g_{NKY} coupling constant [16]. Therefore, in this work we will also employ the $\gamma_5\sigma_{\mu\nu}$ structure to evaluate the $ND\Lambda_c$ form factor and coupling constant, and compare our results with the previous evaluation [11] carried out in the $i\not{q}\gamma_5$ structure. In the case of $ND^*\Lambda_c$ vertex, we will study the sum rule based on the $\not{p}'\gamma_\mu\not{p}$ structure since, in general, the structures with a large number of γ matrices are more likely to be stable.

The phenomenological side of the vertex function is obtained by the consideration of the Λ_c and N intermediate states contribution to the matrix element in Eq.(1):

$$A_M^{(phen)}(p,p',q) = \lambda_{\Lambda_c}\lambda_N \frac{(\not{p}' + M_{\Lambda_c})}{p'^2 - M_{\Lambda_c}^2} g_M(q^2) \frac{(\not{p} + M_N)}{p^2 - M_N^2} + \text{higher resonances} \ , \tag{6}$$

where λ_{Λ_c} and λ_N are the couplings of the currents with the respective hadronic states and the meson couplings are given by

$$g_5(q^2) = i\gamma_5 \frac{m_D^2 f_D}{m_c} \frac{g_{ND\Lambda_c}(q^2)}{q^2 - m_D^2} \ , \tag{7}$$

$$g_\mu(q^2) = m_{D^*} f_{D^*} \frac{g_{ND^*\Lambda_c}(q^2)}{q^2 - m_{D^*}^2} \left(-\gamma_\mu + \frac{\not{q} q_\mu}{m_{D^*}^2} \right) \ , \tag{8}$$

where m_D, m_D^*, f_D and f_{D^*} are the mass and decay constant of the mesons D and D^* respectively, and m_c is the c quark mass.

We will write a sum rule in the structure $\sigma^{\mu\nu}\gamma_5 p'_\mu p_\nu$ for $g_{ND\Lambda_c}(q^2)$ and one in the structure $\not{p}'\gamma_\mu\not{p}$ for $g_{ND^*\Lambda_c}(q^2)$ and we call F the invariant amplitude associated with these structures. The contribution of higher resonances and continuum in Eq. (6) will be taken into account as usual in the standard form of ref. [17].

In the OPE side only even dimension operators contribute to the chosen structures, since the dimension of Eq.(1) is four and $p'p$ takes away two dimensions. The diagrams that contribute, after a double Borel transformation, up to dimension six are shown in Fig. 1. The gluon condensate also contributes, but it always appears with a large suppression factor which arises from the two-loop internal momentum integration. Therefore, its contribution is of little influence and will be neglected. To evaluate the perturbative contribution (Fig. 1a) we write a double dispersion relation to the invariant amplitude, F, over the virtualities p^2 and p'^2 holding $Q^2 = -q^2$ fixed, and use the Cutkosky's rules [18] to evaluate the double discontinuity (see ref. [17]). After doing a double Borel transformation [17] in both variables $P^2 = -p^2 \rightarrow M^2$ and $P'^2 = -p'^2 \rightarrow M'^2$, and subtracting the continuum contribution, we get

$$\left[\tilde{F}(M^2, M'^2, Q^2)\right]_a = -\frac{1}{4\pi^2} \int_{m_c^2}^{u_0} du \int_0^{s_0} ds\, \rho(u,s,Q^2)\, e^{-u/M'^2} e^{-s/M^2} \ , \tag{9}$$

with

$$\rho(u,s,Q^2) = \pm\frac{3}{8\pi^2}\frac{1}{\sqrt{\lambda(s,u,Q^2)}}\int_0^s dm^2\ \left\{m^2\left(-1+\frac{m_c^2(s-u-Q^2)+(s-u)^2+Q^2(s+u)}{\lambda(s,u,Q^2)}\right)\right.$$
$$\left. +\ \frac{2m^4Q^2}{\lambda(s,u,Q^2)}\right\}\Theta(1-(\overline{\cos\theta_K})^2)\Theta\left(u-Q^2-m_c^2+\frac{Q^2u}{m_c^2}-s\right)\ , \tag{10}$$

where

$$\overline{\cos\theta_K} = 2s\frac{u+m^2-m_c^2-p_0'(s+m^2)/\sqrt{s}}{(s-m^2)\sqrt{\lambda(s,u,Q^2)}}\ , \tag{11}$$

with $p_0' = (s+u+Q^2)/(2\sqrt{s})$ and $\lambda(s,u,Q^2) = s^2+u^2+Q^4-2su+2Q^2s+2Q^2u$.

In Eq.(9) $\tilde{F}$ stands for the double Borel transformation of the amplitude F, and the subscript a refers to the diagram in Fig. 1a. u_0 and s_0 give the continuum thresholds for the baryons Λ_c and nucleon respectively, which are, in general, taken from the mass sum rules.

The next lowest dimension operator is the quark c mass times the quark condensate with dimension four (Fig. 1b). Since we are neglecting the light quark masses, only terms proportional to $m_c\langle\overline{q}q\rangle$ will appear. These terms give, after the double Borel transformation

$$\left[\tilde{F}(M^2,{M'}^2,Q^2)\right]_b = -\frac{m_c\langle\overline{q}q\rangle}{4\pi^2}\int_{m_c^2}^{u_0}du\int_0^{s_0}ds\,\alpha(s,u,Q^2)e^{-u/{M'}^2}e^{-s/M^2}\ , \tag{12}$$

where

$$\alpha(s,u,Q^2) = \pm\frac{s(2m_c^2+s-u+Q^2)}{(\lambda(s,u,Q^2))^{3/2}}\Theta\left(u-Q^2-m_c^2+\frac{Q^2u}{m_c^2}-s\right)\ . \tag{13}$$

The next contribution comes from the diagram involving dimension 6 operator of the type $\langle\overline{q}q\overline{q}q\rangle$ ($\simeq\langle\overline{q}q\rangle^2$) shown in Fig. 1c:

$$\left[\tilde{F}(M^2,{M'}^2,Q^2)\right]_c = \pm\frac{\langle\overline{q}q\rangle^2}{3}e^{-m_c^2/{M'}^2}\ . \tag{14}$$

In Eqs. (10) (13) and (14) the $+$ and $-$ sign refers to the $ND\Lambda_c$ and $ND^*\Lambda_c$ vertices repectively.

The Borel transformation of the phenomenological side gives

$$\left[\tilde{F}(M^2,{M'}^2,Q^2)\right]_{phen} = \lambda_{\Lambda_c}\lambda_N G_{NM\Lambda_c}(Q^2)e^{-M_N^2/M^2}e^{-M_{\Lambda_c}^2/{M'}^2}\ , \tag{15}$$

where the continuum contribution has already been incorporated in the OPE side, through the continuum thresholds s_0 nd u_0. In Eq. (15) we have defined

$$G_{NM\Lambda_c}(Q^2) = \frac{m_D^2 f_D}{m_c}\frac{g_{ND\Lambda_c}(Q^2)}{Q^2+m_D^2}\quad\text{for } M=D$$
$$= m_{D^*}f_{D^*}\frac{g_{ND^*\Lambda_c}(Q^2)}{Q^2+m_{D^*}^2}\quad\text{for } M=D^*\ . \tag{16}$$

In order to obtain $g_{ND(D^*)\Lambda_c}(Q^2)$ we identify Eq. (15) with the sum of Eqs. (9), (12) and (14). We obtain:

$$g_{ND\Lambda_c}(Q^2) = \frac{e^{M_N^2/M^2} e^{M_{\Lambda_c}^2/M'^2}}{\lambda_{\Lambda_c}\lambda_N} \frac{m_c(Q^2+m_D^2)}{m_D^2 f_D} \left[-\frac{1}{4\pi^2}\int_{m_c^2}^{u_0} du \int_0^{s_0} ds\, e^{-u/M'^2} e^{-s/M^2} \Big(\rho(u,s,Q^2) \right.$$
$$\left. +\; m_c\langle\overline{q}q\rangle\, \alpha(s,u,Q^2)\Big) + \frac{\langle\overline{q}q\rangle^2}{3} e^{-m_c^2/M'^2} \right] , \tag{17}$$

and

$$g_{ND^*\Lambda_c}(Q^2) = -\frac{e^{M_N^2/M^2} e^{M_{\Lambda_c}^2/M'^2}}{\lambda_{\Lambda_c}\lambda_N} \frac{Q^2+m_{D^*}^2}{m_{D^*} f_{D^*}} \left[-\frac{1}{4\pi^2}\int_{m_c^2}^{u_0} du \int_0^{s_0} ds\, e^{-u/M'^2} e^{-s/M^2} \Big(\rho(u,s,Q^2) \right.$$
$$\left. +\; m_c\langle\overline{q}q\rangle\, \alpha(s,u,Q^2)\Big) + \frac{\langle\overline{q}q\rangle^2}{3} e^{-m_c^2/M'^2} \right] , \tag{18}$$

For λ_{Λ_c} and λ_N we use the expressions obtained from the respective mass sum rules for Λ_c [13] and for the nucleon [12,19]:

$$|\lambda_{\Lambda_c}|^2 = e^{M_{\Lambda_c}^2/{M'_M}^2} \left\{ \frac{m_c^4}{512\pi^4}\int_{m_c^2}^{u_o} du\, e^{-u/{M'_M}^2} \left[\left(1-\frac{m_c^4}{u^2}\right)\left(1-\frac{8u}{m_c^2}+\frac{u^2}{m_c^4}\right) - 12\ln\left(\frac{m_c^2}{u}\right)\right] \right.$$
$$\left. +\; \frac{\langle\overline{q}q\rangle^2}{6} e^{-m_c^2/{M'_M}^2} \right\} , \tag{19}$$

$$|\lambda_N|^2 = e^{M_N^2/M_M^2}\left(\frac{M_M^6}{32\pi^4}E_2 + \frac{2}{3}\langle\overline{q}q\rangle^2\right) , \tag{20}$$

where $E_2 = 1 - e^{-s_0/M_M^2}(1 + s_0/M_M^2 + s_0^2/(2M_M^4))$ accounts for the continuum contribution. For consistency we have also neglected the contribution of the gluon condensate in the mass sum rules, since it is of little influence.

In Eqs. (19) and (20) M_M^2 and ${M'_M}^2$ represent the Borel masses in the two-point function of the nucleon and Λ_c respectively. Comparing Eqs. (17), (19) and (20) we can see that the exponentials multiplying Eq. (17) disappear if we choose

$$2M_M^2 = M^2 \qquad \text{and} \qquad 2{M'_M}^2 = M'^2 . \tag{21}$$

Indeed, this way of relating the Borel parameters in the two- and three-point functions is a crucial ingredient for the incorporation of heavy quark symmetries, and leads to a considerable reduction of the sensitivity to input parameters, such as continuum thresholds s_0 and u_0, and to radiative corrections [14,20].

The parameter values used in all calculations are $m_c = 1.5$ GeV, $m_D = 1.87$ GeV, $m_{D^*} = 2.01$ GeV,$M_N = 938$ MeV, $M_{\Lambda_c} = 2.285$ GeV, $f_D = 170$ MeV $f_{D^*} = 240$ MeV [21], $\langle\overline{q}q\rangle = -(0.23)^3$ GeV^3. We parametrize the continuum thresholds as

$$s_0 = (M_N + \Delta_s)^2 , \tag{22}$$

and

$$u_0 = (M_{\Lambda_c} + \Delta_u)^2 \; . \tag{23}$$

The values of u_0 and s_0 are extracted from the two-point function sum rules for M_{Λ_c} and M_N in Eqs. (19) and (20) and the respective sum rules in the **1** structure given in refs. [12,13,19]. We found a good stability for M_{Λ_c} and M_N, being able to reproduce the experimental values for the masses in the Borel mass $M_M^2 \sim 1\,\mathrm{GeV}^2$ and ${M'_M}^2 \sim 6\,\mathrm{GeV}^2$, with $\Delta_s = 0.7\,\mathrm{GeV}$ and $\Delta_u = 0.6\,\mathrm{GeV}$, which are the values that we are going to use in the calculations.

To allow for different values of M^2 and M'^2 we take them proportional to the respective baryon masses. In this way we study the sum rule as a function of M^2 at a fixed ratio

$$\frac{M^2}{M'^2} = \frac{M_N^2}{M_{\Lambda_c}^2} \; . \tag{24}$$

In Fig. 2 we show the behavior of the perturbative, quark condensate and four quark condensate contributions to the form factor $g_{ND\Lambda_c}(Q^2)$ at $Q^2 = 0$ GeV^2 (since in this case the cut in the t channel starts at $t \sim m_c^2$ and thus the Euclidian region stretches up to that threshold) as a function of the Borel mass M^2. We observe that the different contributions add to give a very stable result as a function of the Borel mass. Since the Borel masses in the two- and three-point functions are related by Eq. (21) and since $M_M^2 \sim 1\,\mathrm{GeV}^2$, to study the Q^2 dependence of the form factor we fix $M^2 = 2.5\,\mathrm{GeV}^2$ where the perturbative contribution is the dominant one. The behaviour of the curve for other Q^2 values is similar, however, for $Q^2 > 1.5\,\mathrm{GeV}^2$ the perturbative contribution is no longer the dominant one. Therefore, in Fig. 3 we show the Q^2 dependence of the form factor in the region $0 \leq Q^2 \leq 1.5\,\mathrm{GeV}^2$ (dots) where we believe we can trust the QCDSR results. We see that the form factor is practically constant in this region, showing a tiny decrease. We can fit the QCDSR results with a monopole form, as can be seen by the solid line in Fig. 3, and we get $g_{ND\Lambda_c}(Q^2) = 843/(102+Q^2)$ which corresponds to a cut-off of order of 10 GeV, much bigger than the values used in ref. [8]. If we vary the value of the Borel mass used to extract the Q^2 dependence of the form factor we can even get a result that shows a tiny increase in the considered region. Therefore, in view of the uncertainties involved in the approach, we can say that the QCDSR give a constant value for the form factor. Considering 20% variation in the continuum thresholds and f_D varying in the interval $f_D = 170 \pm 10$ MeV [21] we get:

$$g_{ND\Lambda_c}(Q^2) = 7.9 \pm 0.9 \; , \tag{25}$$

in agreement with our previous estimate of the coupling constant [11].

The same analysis can be done for $g_{ND^*\Lambda_c}(Q^2)$ and in Fig. 4 we show the perturbative, quark condensate and four quark condensate contributions to the form factor $g_{ND^*\Lambda_c}(Q^2)$ at $Q^2 = 0$ GeV^2 as a function of the Borel mass M^2. Since the OPE sides of both form factors differ only by a sign the aspect of Fig. 2 is very similar to Fig. 4. In Fig. 5 we show the Q^2 dependence of the form factor $g_{ND^*\Lambda_c}(Q^2)$, extracted at $M^2 = 2.5\,\mathrm{GeV}$, in the region $0 \leq Q^2 \leq 1.5\,\mathrm{GeV}^2$ (dots). Again we observe that the form factor is practically constant in this region, showing a tiny decrease. Fitting the QCDSR results with a monopole form (solid line in Fig. 5) we get $g_{ND^*\Lambda_c}(Q^2) = -204/(26.2 + Q^2)$ which corresponds to a cut-off

of order of 5 GeV, still much bigger than the values used in ref. [8]. Despite the fact that for this particular choice of M^2, s_0 and u_0 we get a form factor with a smaller cut-off than obtained for $g_{ND\Lambda_c}(Q^2)$, varying M^2 and the continuun thresholds we still can get results that show a tiny increase in the considered region. Therefore, also in this case, due to the uncertainties involved in the approach, we conclude that the QCDSR give a constant value for the form factor. Considering 20% variation in the continuum thresholds and f_{D^*} varying in the interval $f_{D^*} = 240 \pm 20$ MeV [21] we get:

$$g_{ND^*\Lambda_c}(Q^2) = -7.5 \pm 1.1 \,, \tag{26}$$

It is important to mention that constant form factors in hadronic vertices have already been found in the context of QCDSR. This is the case of the $g_{DD^*\pi}$ form factor. This form factor has been studied in the QCDSR approach for an off-shell D^* [21–23], and for an off-shell pion [24]. As a matter of fact, in refs. [21–23] the authors analyze the semileptonic $f_+(t)$ form factor defined by

$$\langle \pi(p_\pi)|\bar{u}\gamma_\mu c|D(p_D)\rangle = (p_\pi + p_D)_\mu f_+(t) + (p_D - p_\pi)_\mu f_-(t) \,, \tag{27}$$

where $t = q^2 = (p_D - p_\pi)^2$. Since the vector current $V_\mu = \bar{u}\gamma_\mu c$ has the same quantum numbers as the vector meson D^*, the same sum rules studied in refs. [21–23] can be used to study the hadronic form factor $g_{DD^*\pi}(t)$, for an off-shell D^* meson. It is only the phenomenological side of the sum rule that has to be modified to allow for a coupling between the vector current and the vector meson:

$$\begin{aligned}\langle \pi(p_\pi)|V_\mu|D(p_D)\rangle &= \langle 0|V_\mu|\pi(-p_\pi)D(p_D)\rangle \\ &= \langle 0|V_\mu|D^*(q)\rangle \frac{1}{t - m_{D^*}^2} \langle D^*(q)\pi(p_\pi)|D(p_D)\rangle \,. \end{aligned} \tag{28}$$

The $g_{DD^*\pi}(t)$ form factor is defined by the strong amplitude

$$\langle D^*(q,\epsilon)\pi(p_\pi)|D(p_D)\rangle = g_{DD^*\pi}(p_D + p_\pi)^\alpha \epsilon_\alpha \,. \tag{29}$$

Therefore, using the the vacuum to vector meson transition amplitude defined in terms of the vector meson decay constant f_{D^*}:

$$\langle D^*(q,\epsilon)|V_\mu|0\rangle = m_{D^*} f_{D^*} \epsilon_\mu^* \,; \tag{30}$$

we can rewrite Eq. (28) as

$$\langle \pi(p_\pi)|V_\mu|D(p_D)\rangle = -\frac{m_{D^*} f_{D^*} g_{DD^*\pi}(t)}{t - m_{D^*}^2} \left((p_\pi + p_D)_\mu - \frac{q_\mu (p_D^2 - p_\pi^2)}{m_{D^*}^2} \right) \,. \tag{31}$$

Comparing Eqs. (27) and (31) we imediately see that $g_{DD^*\pi}(t)$ and $f_+(t)$ are related by

$$f_+(t) = \frac{m_{D^*} f_{D^*} g_{DD^*\pi}(t)}{m_{D^*}^2 - t} \,. \tag{32}$$

The authors of refs. [21,22] claim that the QCDSR results for $f_+(t)$ can be well fitted by a monopole form with a pole mass $m_{pol} = m_{D^*}$. According with Eq. (32) this implies a constant $g_{DD^*\pi}(t)$. More recently the authors of [23], tried to fit their QCDSR results with a double pole parametrization of the type

$$f_+(t) = \frac{f_+(0)}{(1 - t/m_{D^*}^2)(1 - \alpha t/m_{D^*}^2)} \; . \tag{33}$$

They have obtained $\alpha = 0.01 \begin{array}{c} +0.11 \\ -0.07 \end{array}$ and, therefore, they have concluded that α is consistent with zero, which means the complete dominance of the vector meson pole to f_+ with a consequent prediction of a constant hadronic form factor in the $DD^*\pi$ vertex.

It is interesting to notice that, due to the uncertainties in the value of α given above, one can not even rule out a $DD^*\pi$ form factor that slightly growns in the Euclidian region, as mentioned above in the case of $ND(D^*)\Lambda_c$.

A very different result for the form factor in the $DD^*\pi$ vertex was obtained in the case of an off-shell pion [24]. This form factor is important in the processes described in ref. [25], which could contribute to explain the enhacement in the production of dileptons of intermediate masses, also observed in the NA50 experiment. In the case that the pion is off-shell, the form factor shows a very pronounced Q^2 dependence, with Q^2 being the squared momentum of the off-shell pion in the Euclidian region. This result might be suggesting that the size of a hadronic vertex depends on which particle is off-shell. If the off-shell particle is light, then the vertex is not point like. However, if the off-shell particle is heavy, then the vertex is point like with a consequent constant form factor. Our result for the $ND(D^*)\Lambda_c$ form factors corroborate this hypothesis.

In conclusion, in this work we have calculated the form factors for the hadronic vertices $ND\Lambda_c$ and $ND^*\Lambda_c$ using QCD sum rules. These form factors are important to evaluate the charmonium dissociation cross section by hadrons, in the framework of meson exchange models based on effective lagrangians. In the construction of the sum rules we have performed a double Borel transformation with respect to the nucleon and Λ_c momenta, and we have evaluated the form factors as a function of the heavy meson momentum Q^2. We have studied the sum rules in the structures $\sigma^{\mu\nu}\gamma_5 p'_\mu p_\nu$ for $g_{ND\Lambda_c}(Q^2)$ and $\not{p}'\gamma_\mu\not{p}$ for $g_{ND^*\Lambda_c}(Q^2)$. In the studied Q^2 region, our results are compatible with constant form factors in these vertices. Considering 20% of variation in the continuum thresholds, and around 10% of variation in the meson decay constants we got:

$$\begin{aligned} g_{ND\Lambda_c}(Q^2) &= 7.9 \pm 0.9 \\ g_{ND^*\Lambda_c}(Q^2) &= -7.5 \pm 1.1 \; . \end{aligned} \tag{34}$$

It is important to stress that constant form factors where also found, in the framework of QCDSR, in the vertex $D^*D\pi$ with an off shell D^*. Off course for very large values of Q^2 the form factors should go to zero, but, according to our results, this means that the values of the cut-offs in these form factors should be much larger than 2 GeV, which is the typical value used in refs. [7–9].

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REFERENCES

[1] NA50 Collaboration (M.C. Abreu et al.), *Phys. Lett.* **B477**, 28 (2000).

[2] W. Cassing, E.L. Bratkovskaya and S. Juchem, nucl-th/0001024, *Nucl. Phys.* **A674**, 249 (2000).

[3] A. Capella, E.G. Ferreiro and A.B. Kaidalov, hep-ph/0002300, *Phys. Rev. Lett.* **85**, 2080 (2000).

[4] A. Sibirtsev, K. Tsushima, K. Saito and A.W. Thomas, nucl-th/9904015, *Phys. Lett.* **B484**, 23 (2000).

[5] S.C. Matinyan and B. Müller, nucl-th/9806027, *Phys. Rev.* **C58**, 2994 (1998).

[6] K.L. Haglin, nucl-th/9907034, *Phys. Rev.* **C61**, 031902 (2000).

[7] Z. Lin and C.M. Ko, nucl-th/9912046, *Phys. Rev.* **C62**, 034903 (2000).

[8] A. Sibirtsev, K. Tsushima and A.W. Thomas, nucl-th/0005041.

[9] K.L. Haglin and C. Gale, nucl-th/0010017.

[10] M.A. Shifman, A.I. Vainshtein and V.I. Zakharov, *Nucl. Phys.* **B120**, 316 (1977).

[11] F.S. Navarra and M. Nielsen, *Phys. Lett.* **B443**, 285 (1998).

[12] B.L. Ioffe, *Nucl. Phys.* **B188**, 317 (1981).

[13] E. Bagan, M. Chabab, H.G. Dosch and S. Narison, *Phys. Lett.* **B278**, 367 (1992); **B287**, 176 (1992); **B301**, 243 (1993).

[14] R.S. Marques de Carvalho et al., *Phys. Rev.* **D60**, 034009 (1999).

[15] H. Kim, S.H. Lee and M. Oka, *Phys. Lett.* **B453**, 199 (1999); *Phys. Rev.* **D60**, 034007 (1999); H. Kim, T. Doi, M. Oka and S.H. Lee, nucl-th/0002011.

[16] M.E. Bracco, F.S. Navarra and M. Nielsen, *Phys. Lett.* **B454**, 346 (1999).

[17] B.L. Ioffe and A.V. Smilga, *Nucl. Phys.* **B216** 373 (1983); *Phys. Lett.* **B114**, 353 (1982).

[18] R.E. Cutkosky, *J. Math Phys.* **1**, 429 (1960).

[19] L.J. Reinders, H. Rubinstein and S. Yazaki, *Phys. Rep.* **127**, 1 (1985).

[20] E. Bagan, P. Ball and P. Gosdzinsky, *Phys. Lett.* **B301**, 249 (1993).

[21] V.M. Belyaev, V.M. Braun, A. Khodjamirian and R. Rückl, *Phys. Rev.* **D51**, 6177 (1995).

[22] P. Ball, V.M. Braun and H.G. Dosch, *Phys. Lett.* **B273**, 316 (1991).

[23] A. Khodjamirian et al., hep-ph/0001297, to appear in *Phys. Rev. D.*

[24] F.S. Navarra, M. Nielsen, M.E. Bracco, M. Chiapparini and C.L. Schat, hep-ph/0005026, *Phys. Lett.* **B489**, 319 (2000).

[25] Z. Lin, C.M. Ko and B. Zhang, *Phys. Rev.* **C61**, 024904 (2000).

FIGURES

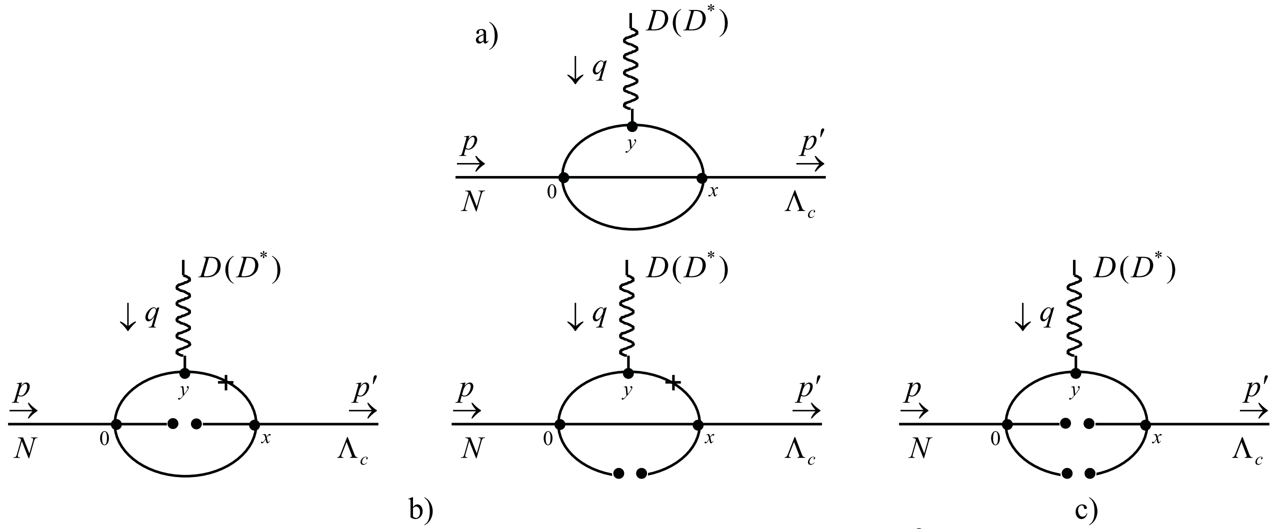

FIG. 1. Diagrams that contribute to the OPE side of $g_{ND(D^*)\Lambda_c}(Q^2)$. The cross stands for a charm mass insertion.

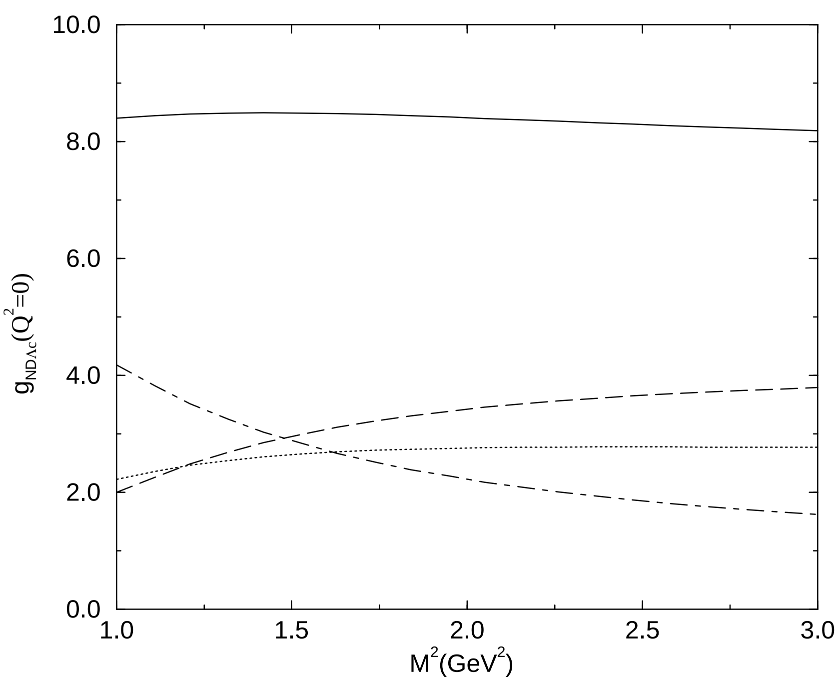

FIG. 2. Borel dependence of the perturbative (dashed line), quark condensate (dotted line) and four quark condensate (dot-dashed line) contributions to the $ND\Lambda_c$ form factor (solid line) at $Q^2 = 0$.

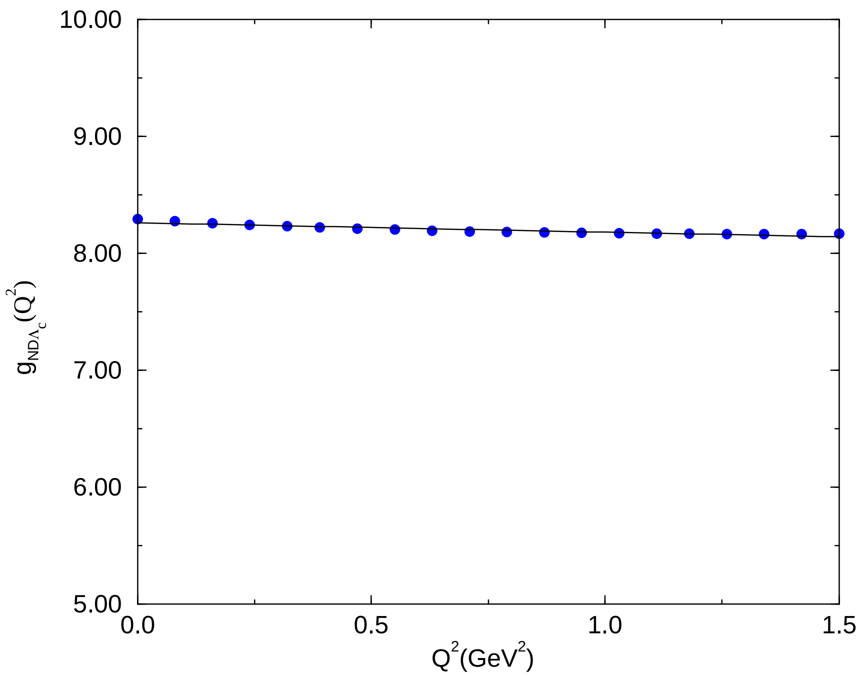

FIG. 3. Momentum dependence of the $ND\Lambda_c$ form factor (dots). The solid line give the parametrization of the QCDSR results with a monopole form.

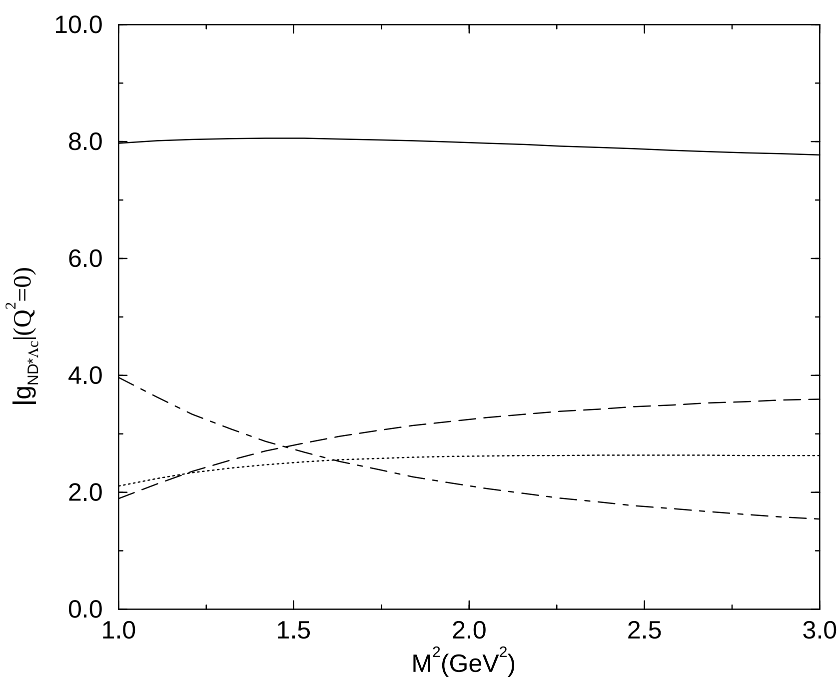

FIG. 4. Borel dependence of the perturbative (dashed line), quark condensate (dotted line) and four quark condensate (dot-dashed line) contributions to the $ND^*\Lambda_c$ form factor (solid line) at $Q^2 = 0$.

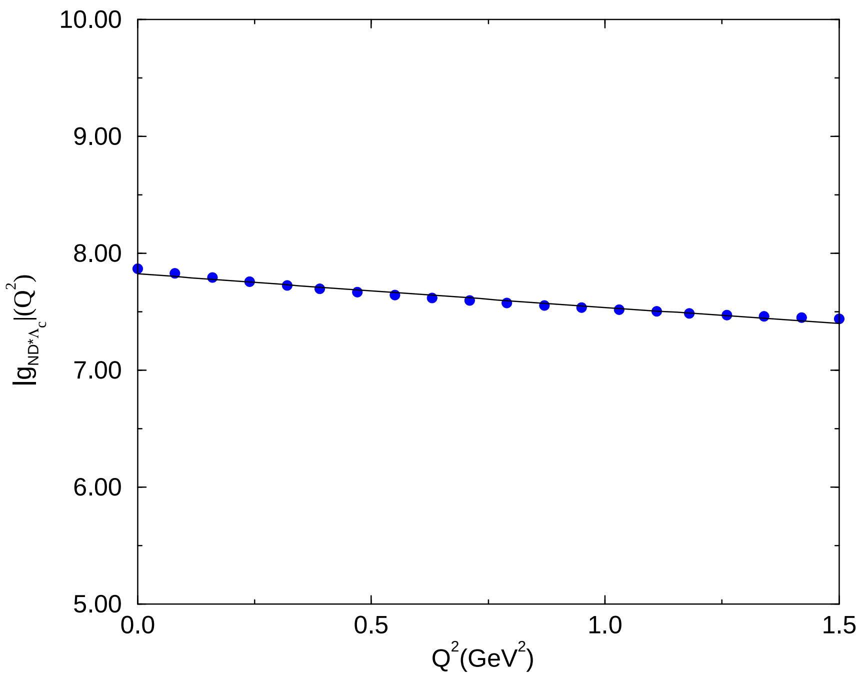

FIG. 5. Momentum dependence of the $ND^*\Lambda_c$ form factor (dots). The solid line give the parametrization of the QCDSR results with a monopole form.